Fermi National Accelerator Laboratory

Fermilab-Pub-89/256-T
BUHEP-89-37
December 14, 1989

Observing the Techniomega at the SSC

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If electroweak symmetry is broken by technicolor, there will be technimesons similar to the mesons of the ordinary strong interactions. In this work, we consider the production and observation of the techniomega (the analogue of the $\omega(783)$) at the SSC.

Operated by Universities Research Association Inc. under contract with the United States Department of Energy

Technicolor theories [1] provide an interesting mechanism for breaking the $SU(2) \times U(1)$ electroweak symmetry down to the $U(1)$ of electromagnetism. These theories predict the existence of meson- and baryon-like particles, the observation of which would indicate that a fermion condensate is responsible for electroweak symmetry breaking.

All of these theories have at least one weak doublet of technifermions. In the simplest models these transform as a fundamental under the new $SU(N_T)$ technicolor (TC) interactions. Except for the gauging of the weak $SU(2) \times U(1)$, this theory respects a chiral $SU(2)_L \times SU(2)_R$, of which the weak generators form a subgroup. When the TC interactions become strong, the left-handed doublet condenses with the right-handed, breaking the chiral group down to its vector component, and simultaneously breaking the weak group to electromagnetism.

In QCD the spontaneous breakdown of the chiral symmetry produces Goldstone bosons, the pions. In technicolor their analogues become the degrees of freedom "eaten" by the W and Z. The other technihadrons should be observable as real particles. In this paper we are concerned with a technivector meson, the analogue of the $\omega(783)$[2].

In the simplest technicolor model, in addition to the usual quarks and leptons, there is one left-handed weak doublet of technifermions (p_L, m_L) and two right-handed singlets, p_R and m_R. The left-handed doublet carries no hypercharge, while p_R and m_R carry hypercharge $+1/2$ and $-1/2$ respectively. Non-minimal technicolor models contain additional technifermions, and their spectrums therefore contain more technihadrons. Though in this work we consider the minimal technicolor model as an example, it is possible to make similar analyses in other models.

The spectrum of the simple theory may be deduced in analogy with QCD. Aside from the three technipions, the "swallowed" degrees of freedom that become the longitudinal components of the W and Z, the spectrum of heavy technimesons begins with the spin-one particles. These are the ρ_T and ω_T, analogues of the $\rho(770)$ and $\omega(783)$. They are respectively a triplet and a singlet under techni-isospin, the vectorial $SU(2)$ which acts on p and m. In addition to these, there will be excited states of the mesons, as well as a full spectrum of technibaryons, the analogues of the nucleons and the Δ's.

We may understand the decays of these particles by using the "equivalence theorem", which states that any amplitude involving one or more longitudinal gauge bosons is equivalent to the same amplitude with the external W_L or Z_L replaced by the appropriate goldstone boson, up to corrections which vanish at high energies[3]. Therefore, in a technicolor theory, the longitudinal components of the W and Z are strongly interacting. For example, in QCD we see the decay $\rho^+ \to \pi^0\pi^+$, and in technicolor we would expect to see $\rho_T^+ \to Z_L W_L^+$.

Since the techniomega is a singlet under techni-isospin, it will not, like the ρ_T, have a TC coupling into a pair of technipions. Therefore, it will not be produced by the gauge boson fusion mechanism[4]. Moreover, in the model considered here, the coupling of the Z to technifermions is purely techni-isotriplet, and so the ω_T is not produced[1] via mixing with the Z. The only allowed TC coupling is to three gauge bosons. Hence, the ω_T is far narrower than the ρ_T, and if it is produced in sufficient quantities it may be easier to extract from the backgrounds.

If the techniomega is not made by the TC interactions, how can one hope for its copious production? The answer is that the TC interactions cannot be the whole story. In any viable model, it is necessary to give the quarks and leptons their observed masses. The quark and lepton masses break electroweak gauge symmetry, so their existence must be related to the formation of the technifermion condensate. The quarks and leptons are usually given a coupling to the technifermions under an extended technicolor (ETC) gauge group. The ETC gauge bosons mediate a coupling of the quarks and leptons into the technifermion condensate. The ETC group is understood to be broken at an ultraheavy scale. At low energies, well below the mass of the ETC gauge bosons, these interactions may be represented as effective four-fermi operators involving fermions and technifermions. The operators may have several different flavor and Lorentz structures, depending on the details of the ETC sector. Consider the operator

$$\frac{1}{4f^2} J_T^\mu J_{f\mu} \tag{1}$$

with

$$J_f^\mu = \left(\bar{u}\gamma^\mu \frac{1-\gamma_5}{2} u + \bar{d}\gamma^\mu \frac{1-\gamma_5}{2} d + \ldots \right) ,$$

[1]In reference [5] the production of the technirho via mixing with the W and Z was computed.

where ... represents all other quarks and leptons, and

$$J_T^\mu = \left(\bar{p}\gamma^\mu \frac{1-\gamma_5}{2} p + \bar{m}\gamma^\mu \frac{1-\gamma_5}{2} m\right) \ .$$

This operator appears naturally in models in which the ETC interactions are approximately flavor symmetric [6], and it is not in contradiction with any current experimental data as long as $f > 1$ TeV. Since it is a singlet under techni-isospin, this operator can be responsible for the production of techniomegas in hadronic supercolliders.

In order to estimate the rate at which this operator produces techniomegas, we use large N_T arguments[7]. In usual large N QCD, one fixes Λ_{QCD} and f_π varies with N. In technicolor theories, we know that v, the analogue of f_π, is about 1/4 TeV. Therefore, we fix v and allow Λ_T to vary with N_T. One finds that

$$M_{\omega T} \approx m_\omega \frac{v}{f_\pi}\left(\frac{3}{N_T}\right)^{\frac{1}{2}} \approx 2069 \text{ GeV} \left(\frac{3}{N_T}\right)^{\frac{1}{2}} \ . \tag{2}$$

Since, to a good approximation, there are no technifermions in the proton and no quarks in the techniomega, we may factorize the operator of equation 1:

$$\langle \omega_T X | \frac{1}{4f^2} J_T^\mu J_{q\mu} | p\bar{p}\rangle = \frac{1}{4f^2}\langle \omega_T | J_T^\mu | 0\rangle \langle X | J_{q\mu} | p\bar{p}\rangle \ , \tag{3}$$

where $|X\rangle$ is the hadronic final state. We may now define $f_{\omega T}$ by

$$\langle \omega_T | J_T^\mu | 0\rangle = \frac{M_{\omega T}^2}{f_{\omega T}} \epsilon^\mu \ . \tag{4}$$

where ϵ^μ is the polarization vector of the techniomega. Here only the vectorial part of J_T^μ contributes. f_ω is defined in analogy to equation 4, but with the $\omega(783)$ replacing ω_T. Then, scaling from QCD, we find

$$f_{\omega_T} = f_\omega \left(\frac{3}{N_T}\right)^{\frac{1}{2}} \ . \tag{5}$$

The determination of f_ω may be made in several ways. Naively, it would appear that the most direct method is to use the decay $\omega \to e^+e^-$. The partial width is $4\pi\alpha m_\omega/(27 f_\omega^2)$, yielding $f_\omega = 5.7$. Unfortunately the situation the $\phi(1020)$, the isosinglet vector meson which is mostly $s\bar{s}$. One expects that f_ϕ is affected by the mass of the strange quark, an effect which will be absent in the technicolor case, since bare technifermion masses would violate the gauge symmetry. To avoid this complication one may instead choose to use the approximate nonet symmetry, and assume that $f_\omega = f_\rho$. The partial width of the decay $\rho \to e^+e^-$ is $4\pi\alpha m_\omega/(3f_\rho^2)$, and so we find $f_\rho = 5.0$; this is the value for f_ω we use throughout. There is evidently about a 20% uncertainty in the value of f_ω.

One may now compute the total production of techniomegas in a hadronic collider using the narrow width approximation. The cross section is given by

$$\sigma_{\omega T} = \frac{M_{\omega T}^4 \pi}{96 f_{\omega T}^2 f^4} \int_0^1 dx \left(f_u^{(1)}(x) f_{\bar{u}}^{(2)} \left(\frac{M_{\omega T}^2}{xs} \right) + (u \to d) + ((1) \leftrightarrow (2)) \right) \ . \quad (6)$$

Denoting the integral in this equation by $\mathcal{L}_{q\bar{q}}(M_{\omega T}^2/s)$, using $\sqrt{s} = 40$ TeV, and the values above for $M_{\omega T}$ and $f_{\omega T}$, we find

$$\sigma_{\omega T} = 6.2 \times 10^{-6} \text{nb} \left(\frac{3}{N_T} \right) \left(\frac{1 \text{ TeV}}{f} \right)^4 \mathcal{L}_{q\bar{q}} \left(\frac{M_{\omega T}^2}{s} \right) \ .$$

We evaluate $\mathcal{L}_{q\bar{q}}$ using the EHLQ set II structure functions[5], choosing the scale $Q^2 = \hat{s}$. The result is a substantial number of techniomegas per year. Table 1 shows the annual production of techniomegas for different values of N_T, assuming the SSC has a CM energy of 40 TeV and a luminosity of 10^{40} cm^2/year. Unfortunately, only a small fraction of these particles decay into observable channels.

The leading decay of the $\omega(783)$ is to $\pi^+\pi^-\pi^0$, which proceeds via the strong interactions. Since the pions are goldstone bosons, we can deduce that the amplitude for this decay must be of the form

$$a(\omega \to \pi^+\pi^-\pi^0) \propto \epsilon^{\mu\nu\rho\sigma} \omega_\mu \partial_\nu \pi^+ \partial_\rho \pi^- \partial_\sigma \pi_0 + O(\partial^5) \ . \quad (7)$$

The terms with higher numbers of derivatives should be small over most of phase space. Since real world pions have mass, the decay is suppressed relative to what it would be if the pions were true goldstone bosons, both by the derivatives in the above matrix element and by phase space. The decay rate is thus reduced by a factor of about 0.245. The technipions, on the other hand, are the longitudinal components of the W and Z, which have a negligible mass relative to the techniomega. Scaling from QCD we expect

$$\begin{aligned}\Gamma(\omega_T \to W_L^+ W_L^- Z_L) &= \left(\frac{3}{N_T}\right)^2 \left(\frac{1}{0.245}\right) \Gamma(\omega \to \pi^+\pi^-\pi^0) \frac{M_{\omega T}}{m_\omega} \\ &\approx 81.6 \text{ GeV} \left(\frac{3}{N_T}\right)^{\frac{5}{2}} , \end{aligned} \tag{8}$$

where we have used the measured value for the partial width[8].

Though this is the leading decay of the ω_T, it probably cannot be observed. The leptonic branching ratios are $BR(W \to \ell\nu) = 22\%$ and $BR(Z \to \ell^+\ell^-) = 6.7\%$, (where $\ell = e, \mu$, and $m_{top} > M_W$), so only 0.1% of the three gauge boson events will decay to purely leptonic states. Moreover, the loss of the two neutrinos means that the invariant mass of the event will not be reconstructed. Therefore, the sharp peak will be smeared out, rendering the techniomega difficult to find above the backgrounds. Recent work has claimed that it may be possible to find the heavy Higgs boson in the channel W^+W^- where one of the W's decays hadronically[9]. If it is possible to identify a W decaying to hadrons in this case, then some of these problems may be obviated. Considerable work would be required to show there is a set of cuts that will pull the signal out of the background of events with two gauge bosons plus two jets.

The next-to-leading decay of the $\omega(783)$ is to $\pi^0\gamma$; it proceeds electromagnetically and has a branching ratio of about 8% [8]. The reason it can compete with the strong decay to three pions is that the latter is suppressed by the three body final state phase space. The analogous mode for the technipion would be to $Z_L\gamma$, which has a very clean signature and a large branching ratio to observable final states. It is this mode which may permit the observation of the ω_T at the SSC.

In order to discuss this decay, we need to construct an effective lagrangian for the interactions of the ω_T with the gauge bosons. We use the techniques of reference [10]. The technipions are exponentiated in a field Σ

$$\Sigma = \exp\left(\frac{i\vec{\pi}_T \cdot \vec{\sigma}}{v}\right) \tag{9}$$

which transforms linearly under the spontaneously broken chiral $SU(2)_L \times$

$SU(2)_R$ of this model. That is

$$\Sigma \to L\Sigma R^\dagger \ . \tag{10}$$

Σ is a 2×2 unitary matrix. We define the field ξ by $\xi\xi = \Sigma$, so that ξ transforms as

$$\xi \to L\xi U^\dagger = U\xi R^\dagger \ . \tag{11}$$

This implicitly defines U as a function of L, R, and the technipion fields. In the case of vectorial techni-isospin transformation $L = R$, we see that $U = L$, independent of the technipions. We next construct the field M^μ, the vector meson field, by

$$M^\mu = \omega_T^\mu I + i\vec{\sigma} \cdot \vec{\rho_T^\mu} \ . \tag{12}$$

We take this to transform as $M^\mu \to UMU^\dagger$, so that ρ_T is a techni-isotriplet and ω_T is a techni-isosinglet[2]. Any other transformation rule is equivalent.

Thus far, the transformations L and R have been understood to be members of a global chiral $SU(2)$. As we know, the $SU(2)_L$ and the third component of the $SU(2)_R$ are gauged. Therefore, we must replace all derivatives in our lagrangian by gauge covariant derivatives. For example, the kinetic energy term for the technipions becomes

$$\mathcal{L}_{KE} = \frac{f^2}{4}(D^\mu\Sigma)^\dagger(D_\mu\Sigma) \ , \tag{13}$$

where

$$D^\mu\Sigma = \partial^\mu\Sigma - igW^\mu \cdot \frac{\vec{\sigma}}{2}\Sigma + ig'\Sigma B^\mu \frac{\sigma_3}{2} \ ,$$

and g and g' are the $SU(2)_L$ and $U(1)_Y$ coupling constants, respectively. By choosing a position dependent gauge transformation $L(x) = \Sigma^\dagger(x)$, $R(x) = I$, we may set $\Sigma = I$, and the technipion fields do not propagate. This is the unitary gauge, in which the W and Z have "swallowed" the technipions. The kinetic energy term of equation 13 becomes a mass term for the gauge bosons, and the relationship $M_W = M_Z \cos(\theta_W)$ is respected.

[2]Strictly speaking, the ρ_T and ω_T need not be assembled into a single representation in this way. By doing so we are imposing a "quartet" symmetry, which would be broken if we wrote down terms in the Lagrangian involving $\mathrm{tr}M^\mu$, which is, after all, invariant under the techni-isospin transformations U.

The term which mediates the weak decay of the vector mesons is

$$\begin{aligned}\mathcal{L}_{Weak} &= \left(\tfrac{N_T}{3}\right)^{\frac{1}{2}} \chi \epsilon_{\mu\nu\rho\sigma} [\quad -g \ \mathrm{tr}(\frac{\vec{\sigma}}{2} \vec{W}^{\rho\sigma} \{\Sigma D^\nu \Sigma^\dagger, \xi M^\mu \xi^\dagger\}) \\ &\qquad + g' \ \mathrm{tr}(\frac{\sigma_3}{2} B^{\rho\sigma} \{\Sigma^\dagger D^\nu \Sigma, \xi^\dagger M^\mu \xi\})] \ , \end{aligned} \tag{14}$$

where $\vec{W}^{\rho\sigma}$ and $B^{\rho\sigma}$ are the field strengths of the $SU(2)$ and $U(1)$ gauge groups respectively. The terms with the anti-commutators replaced by commutators are excluded by CP. The relative sign between the two terms is dictated by the conservation of parity by the TC interactions. We expect that under a parity transformation ($D^\mu \leftrightarrow D_\mu$ and $\Sigma \leftrightarrow \Sigma^\dagger$), and exchange of left and right gauge fields ($B\sigma_3 \leftrightarrow \vec{W} \cdot \vec{\sigma}$), the term goes into itself. One may fit the dimensionless coefficient χ by noting that the decay $\omega(783) \rightarrow \pi^0\gamma$ is described by the analogous lagrangian. The amplitude thus derived is

$$a(\omega \rightarrow \pi^- \gamma) = -8i \frac{\chi e}{f_\pi} \epsilon_{\mu\nu\rho\sigma} \epsilon_\omega^\mu p_\pi^\nu p_\gamma^\rho \epsilon_\gamma^\sigma \ , \tag{15}$$

where ϵ_ω^μ, and ϵ_γ^μ represent the polarization vector of the ω and γ respectively, and e is the electronic charge. Plugging in $\Gamma(\omega \rightarrow \pi^0\gamma) = 680$ keV, we find that $\chi = 2.63 \times 10^{-2}$.

The lagrangian of equation 14 describes all processes in which the ω_T to two gauge bosons, one longitudinal and one transverse. In unitary gauge we find

$$\begin{aligned}\mathcal{L}_{Weak} &= 2ie^2 \quad \left(\frac{N_T}{3}\right)^{\frac{1}{2}} \chi \epsilon^{\mu\nu\rho\sigma} \omega_\mu \Bigg[\left(\frac{1}{\sin^2(\theta_W)}\right) \partial_\rho W_\sigma^+ W_\nu^- \\ &\qquad + \left(\frac{1}{\sin^2(\theta_W)}\right) \partial_\rho W_\sigma^- W_\nu^+ + \left(\frac{\cos^2(\theta_W)}{\sin^2(\theta_W)} + \frac{\sin^2(\theta_W)}{\cos^2(\theta_W)}\right) \partial_\rho Z_\sigma Z_\nu \\ &\qquad + \left(\frac{2}{\sin(\theta_W)\cos(\theta_W)}\right) \partial_\rho \gamma_\sigma Z_\nu \Bigg] \\ &\qquad + \text{terms with more than two gauge bosons} \ , \end{aligned} \tag{16}$$

where θ_W is the weak mixing angle. Using $M_Z = 91.0$ GeV, $M_W = 80.4$ GeV, $\sin^2(\theta_W) = .22$, and $\alpha(m_W) = 1/128$, we derive

$$\Gamma(\omega_T \rightarrow Z\gamma) \ = \ 2.3 \ \text{GeV} \left(\frac{3}{N_T}\right)^{\frac{1}{2}} \ , \tag{17}$$

$$\Gamma(\omega_T \to ZZ) \;=\; 1.05\ \mathrm{GeV}\left(\frac{3}{N_T}\right)^{\frac{1}{2}} , \tag{18}$$

$$\Gamma(\omega_T \to W^+W^-) \;=\; 5.2\ \mathrm{GeV}\ ,\left(\frac{3}{N_T}\right)^{\frac{1}{2}} \tag{19}$$

As in the case of the $\omega(783)$, these weak decays are not negligible.

The techniomega has one further decay, via the same four-fermion operator which produces it. The partial width is

$$\Gamma(\omega_T \to f\bar{f}) = 30\ \mathrm{GeV}\left(\frac{1\ \mathrm{TeV}}{f}\right)^4\left(\frac{5.0}{f_\omega}\right)^2\left(\frac{3}{N_T}\right)^{\frac{3}{2}}\left(\frac{N_f}{24}\right)\ , \tag{20}$$

where N_f is the number of fermion species appearing in J_f^μ. If we believe that all the fermions of the three generations appear, $N_f = 24$. For f about 1 TeV, this decay process produces a considerable number of $\ell^+\ell^-$ events.

Table 2 shows the partial decay widths for the techniomega assuming that $f_\omega = 5$ and N_f=24. One may obtain a rough estimate of the number of observable techniomega events using the entries of table 1 and the partial decay widths given above. However, this would be overly optimistic, since it would include $Z\gamma$ events in which one or both of the leptons from the decay of the Z missed the detector. Instead, we have calculated the complete Feynman diagram for production and decay of the techniomega: $q\bar{q} \to \omega_T \to Z\gamma$. The final state Z is longitudinal, so its decay into $\ell^+\ell^-$ may also be included. In this way, we may make cuts on the actual final state particles. The partial width of the techniomega is included by replacing its propagator by a Breit-Wigner form, using the partial widths given above. The amplitude thus computed is integrated numerically and the desired cuts are made. For the initial state structure functions, EHLQ set II is again used.

Cuts are imposed in order to mimic a realistic experimental situation. We demand that the absolute value of the rapidities of the leptons and the photon be less than 2.5. We also reject events in which the photon or a lepton has transverse momentum less than 50 GeV. The opening angle between the two leptons must be large enough for them to be resolved in a reasonable detector: $\Delta\phi, \Delta y > 0.10$. Lastly, the total invariant mass of the event should be within one and a half techniomega widths of the mass of the techniomega. When the techniomega is narrow, we require instead that the invariant mass be within 50 GeV of the techniomega mass.

The first column of table 3 shows the the number of events per SSC year for various values of N_T and f. As f grows, the production rate shrinks like f^{-4}; this is somewhat mitigated by the shrinkage of the width to $f\bar{f}$, which increases somewhat the branching ratio to $Z\gamma$. However, since the width to three gauge bosons is larger than the width to $f\bar{f}$, this effect is not appreciable. For the values of N_T and f under consideration the numbers of events will go roughly like f^{-4}.

There are several potential backgrounds to this signal. The most obvious is the physics background from continuum $Z\gamma$ production. The second column of table 3 shows the number of $q\bar{q} \to Z\gamma \to \ell^+\ell^-\gamma$ events which pass the cuts above. Another physics background comes from the continuum production of $e^+e^-\gamma$ via an insertion of an ETC generated four-fermion operator like that of equation 1. Such a production mechanism yields far fewer events than continuum $Z\gamma$ production, both because of the large scale f suppressing the process, and also because the lepton pair rather infrequently has an invariant mass near the Z. In addition to these, there are also possible "junk" backgrounds. For example, at the SSC there will be copious production of $Z + jet$ events; a potential problem occurs if the jet is frequently mistaken for a single photon. We estimate that the production of $Z + jets$ is about two orders of magnitude larger than $Z\gamma$, so one requires a jet-photon rejection of only 10^{-3} in order to to suppress this junk background well below the irreducible physics background.

The third column of table 3 shows the numbers of events in which the techniomega decays into an $\ell^+\ell^-$ pair. This process interferes with ordinary Drell-Yan production of leptons, and this column therefore represents the sum of the two processes. We have demanded that the rapidity of the leptons be less than 2.5. This process is suppressed by eight powers of f, four for the production and four for the decay.

We conclude by referring again to the tables. The techniomega is observable in a very clean $Z\gamma$ channel, and it is significantly above the background if f is about 1 TeV. Unfortunately, the production is suppressed by a large power of f, so that it is unlikely that the ω_T can be seen if f is as large as 2 TeV. The situation is better for larger values of N_T, both because the mass of the techniomega is smaller and because the decay width to three gauge bosons or lepton pairs is reduced. Similar conclusions apply to the decay to lepton pairs. There is a small but nonzero window to observe the ω_T at the

SSC.

Acknowledgements

We would like to thank E.H. Simmons for reading the manuscript.

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Table Captions

Table 1: The mass of the techniomega and total number of techniomegas produced in one SSC-year, (10^{40} cm^2 at 40 TeV CM energy) for different values of N_T. The computation was done in the narrow width approximation and includes all techniomegas, irrespective of decay channel.

Table 2: The partial widths of the techniomega into W^+W^-Z, $f\bar{f}$, W^+W^-, ZZ, and $Z\gamma$, for different values of N_T. Here, $f_\omega = 5.0$ and $N_f = 24$.

Table 3: For different values of N_T and f, the three columns show the number of events per SSC-year from three sources: observable $\omega_T \rightarrow Z\gamma \rightarrow e^+e^-$ signal, the continuum $Z\gamma$ background, and the combined rate for $\ell^+\ell^-$ from techniomega decay and Drell-Yan production. See the text for the cuts imposed.

N_T	$M_{\omega T}$	total events/year
2	2530 GeV	18000
4	1790 GeV	49000
6	1460 GeV	83000

N_T	$\Gamma(WWZ)$	$\Gamma(f\bar{f})$	$\Gamma(WW)$	$\Gamma(ZZ)$	$\Gamma(Z\gamma)$	$\mathrm{BR}(Z\gamma)$
2	225 GeV	55 GeV $\left(\frac{1\ \mathrm{TeV}}{f}\right)^4$	6.4 GeV	1.3 GeV	2.8 GeV	1%
4	40 GeV	19 GeV $\left(\frac{1\ \mathrm{TeV}}{f}\right)^4$	4.5 GeV	0.9 GeV	2.0 GeV	3%
6	14 GeV	11 GeV $\left(\frac{1\ \mathrm{TeV}}{f}\right)^4$	3.7 GeV	0.7 GeV	1.6 GeV	5%

N_T		$\omega_T \to \ell^+\ell^-\gamma$	continuum $Z\gamma$	Drell-Yan
2	$f = 1.0$ TeV	9.2	2.2	144
	$f = 1.5$ TeV	2.1	1.8	21
	$f = 2.0$ TeV	0.7	1.8	15
	$f = \infty$	0	1.7	14
4	$f = 1.0$ TeV	48	1.9	346
	$f = 1.5$ TeV	12	1.4	31
	$f = 2.0$ TeV	4.2	1.4	15
	$f = \infty$	0	1.3	13
6	$f = 1.0$ TeV	110	2.0	536
	$f = 1.5$ TeV	30	2.0	51
	$f = 2.0$ TeV	10	2.0	25
	$f = \infty$	0	2.0	22